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# Higgs-Mediated $B^0 \rightarrow \mu^+ \mu^-$ in Minimal Supersymmetry

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## Abstract

In this letter we demonstrate a new source for large flavor-changing neutral currents within the minimal supersymmetric standard model. At moderate to large  $\tan \beta$ , it is no longer possible to diagonalize the masses of the quarks in the same basis as their Yukawa couplings. This generates large flavor-violating couplings of the form  $\bar{b}_R d_L \phi$  and  $\bar{b}_R s_L \phi$  where  $\phi$  is any of the three neutral, physical Higgs bosons. These new couplings lead to rare processes in the  $B$  system such as  $B^0 \rightarrow \mu^+ \mu^-$  decay and  $B^0 - \bar{B}^0$  mixing. We show that the latter is anomalously suppressed, while the former is in the experimentally interesting range. Current limits on  $B^0 \rightarrow \mu^+ \mu^-$  already provide nontrivial constraints on models of moderate to large  $\tan \beta$ , with an observable signal possible at Run II of the Tevatron if  $m_A \lesssim 400 - 700$  GeV, extending to the TeV range if a proposed Run III of  $30 \text{ fb}^{-1}$  were to occur.

Extensions of the Standard Model containing more than one Higgs SU(2) doublet generically allow flavor-violating couplings of the neutral Higgs bosons. Such couplings, if unsuppressed, will lead to large flavor-changing neutral currents, in direct opposition to experiment [1]. Models such as the Minimal Supersymmetric Standard Model (MSSM) avoid these dangerous couplings by segregating the quark and Higgs fields so that one Higgs ( $H_u$ ) can couple only to  $u$ -type quarks while the other ( $H_d$ ) couples only to  $d$ -type. Within unbroken supersymmetry this division is completely natural; in fact, it is required by the holomorphy of the superpotential.

However, after supersymmetry is broken, there is nothing left to protect this division. In fact, it has been known for some time that couplings of the form  $QU^c H_d^*$  and  $QD^c H_u^*$  are generated at one-loop [2]. As such, one would expect some flavor violation to arise in the neutral Higgs sector, but always suppressed by loops and therefore small (as small or smaller than Standard Model flavor-changing). But this is not the correct deduction.

Hall, Rattazzi and Sarid (HRS) [3] showed that at moderate to large  $\tan\beta \equiv \langle H_u \rangle / \langle H_d \rangle$  the contributions to  $d$ -quark masses coming from the non-holomorphic operator  $QD^c H_u^*$  can be equal in size to those coming from the usual holomorphic operator  $QD^c H_d$  despite the loop suppression suffered by the former. This is because the operator itself gets an additional enhancement of  $\tan\beta$ . That is, the product  $\tan\beta/16\pi^2$  need not be very small as  $\tan\beta$  approaches its upper bound of 60 to 70.

The HRS result was followed shortly by Ref. [4] which analyzed the entire  $d$ -quark mass matrix in the presence of these corrections and found appreciable contributions to the CKM mixing angles. It has also recently been realized that the HRS corrections can significantly alter the (flavor-conserving) couplings of the Higgs bosons [5, 6]. In this letter we take our analysis from Ref. [6] one step further and show that flavor-changing couplings of the neutral Higgs bosons are also generated. We will show that these couplings can be appreciable and can be so even without invoking squark mixing and/or non-minimal Kähler potentials [7], and remain large even in the limit of heavy squarks and gauginos. These new couplings generate a variety of flavor-changing processes, including  $\bar{B}^0 - B^0$  mixing and decays such as  $B^0 \rightarrow \mu^+ \mu^-$  which we will study in this letter. A more complete discussion of these and other effects will be found in a forthcoming paper [8].

We begin by writing the effective Lagrangian for the interactions of the two Higgs doublets with the quarks in an arbitrary basis:

$$-\mathcal{L}_{eff} = \bar{D}_R \mathbf{Y}_D Q_L H_d + \bar{D}_R \mathbf{Y}_D [\epsilon_g + \epsilon_u \mathbf{Y}_U^\dagger \mathbf{Y}_U] Q_L H_u^* + h.c. \quad (1)$$

Here  $\mathbf{Y}_D$  and  $\mathbf{Y}_U$  are the  $3 \times 3$  Yukawa matrices of the microscopic theory, while the  $\epsilon_{g,u}$  are the finite, loop-generated non-holomorphic Yukawa coupling coefficients derived by HRS. The leading contributions to  $\epsilon_g$  and  $\epsilon_u$  are generated by the two diagrams in Fig. 1. (There can also be contributions to  $\mathcal{L}_{eff}$  proportional to  $\mathbf{Y}_D \mathbf{Y}_D^\dagger \mathbf{Y}_D$ ; however, since they are typically smaller than the  $\epsilon_g$  contribution and do not generate flavor-violations, we will not consider them further.)

Consider the first diagram in Fig. 1. If all  $\tilde{Q}_i$  masses are assumed degenerate at some scale  $M_{\text{unif}}$  then, at lowest order,  $i = k$  and the diagram contributes only to  $\epsilon_g$ :

$$\epsilon_g \simeq \frac{2\alpha_3}{3\pi} \mu^* M_3 f(M_3^2, m_{\tilde{Q}_L}^2, m_{\tilde{d}_R}^2), \quad (2)$$

where [3]

$$f(x, y, z) = -\frac{xy \log(x/y) + yz \log(y/z) + zx \log(z/x)}{(x-y)(y-z)(z-x)}. \quad (3)$$

Meanwhile, the second diagram of Fig. 1 contributes to  $\epsilon_u$ :

$$\epsilon_u \simeq \frac{1}{16\pi^2} \mu^* A_U f(\mu^2, m_{\tilde{Q}_L}^2, m_{\tilde{u}_R}^2). \quad (4)$$

(We assume that the trilinear  $A$ -terms can be written as some flavor-independent mass times  $\mathbf{Y}_U$ .) For typical inputs, one usually finds  $|\epsilon_g|$  is about 4 times larger than  $|\epsilon_u|$ .

However, there is another sizable contribution to  $\epsilon_u$ , this one coming from the *first* diagram in Fig. 1. It is well-known that  $\tilde{Q}_i$  degeneracy is broken by radiative effects induced by Yukawa couplings. While this would appear to be a higher-order effect, for  $M_{\text{unif}} \gg M_{\text{SUSY}}$  it is amplified by a large logarithm and thus can be  $\mathcal{O}(1)$ . Without resumming that log, one finds a deviation from universality of [9]

$$\Delta \mathbf{m}_{\tilde{Q}}^2 \simeq -\frac{1}{8\pi^2} (3m_0^2 + A_0^2) [\mathbf{Y}_U^\dagger \mathbf{Y}_U + \mathbf{Y}_D^\dagger \mathbf{Y}_D] \log \left( \frac{M_{\text{unif}}}{M_{\text{SUSY}}} \right) \quad (5)$$

where  $m_0$  and  $A_0$  are the common scalar mass and trilinear soft term at  $M_{\text{unif}}$ . (Resummation can bring in additional flavor structure such as  $(\mathbf{Y}_U^\dagger \mathbf{Y}_U)^2$ ,  $(\mathbf{Y}_D^\dagger \mathbf{Y}_D)^2$ ,  $(\mathbf{Y}_U^\dagger \mathbf{Y}_U)(\mathbf{Y}_D^\dagger \mathbf{Y}_D)$ , etc., but these are numerically less significant and do not lead to any new flavor structure.) At the SUSY scale, we can write the  $\tilde{Q}$  mass matrix in the form

$$\mathbf{m}_{\tilde{Q}}^2 = \overline{m}^2 \left( \mathbf{1} + c \mathbf{Y}_U^\dagger \mathbf{Y}_U + c \mathbf{Y}_D^\dagger \mathbf{Y}_D \right) \quad (6)$$

where

$$c \simeq -\frac{1}{8\pi^2} \frac{3m_0^2 + A_0^2}{\overline{m}^2} \log \left( \frac{M_{\text{unif}}}{M_{\text{SUSY}}} \right) \quad (7)$$

and  $\overline{m}^2$  is a flavor-independent mass term. The effect of this non-universality is to generate a contribution to  $\epsilon_u$  proportional to  $\alpha_3$  and thus potentially large (the  $\mathbf{Y}_D^\dagger \mathbf{Y}_D$  piece is again irrelevant). Specifically,

$$\Delta \epsilon_u \simeq \begin{cases} -c\epsilon_g/3 & (m_{\tilde{Q}}^2 \simeq M_3^2) \\ -c\epsilon_g/2 & (m_{\tilde{Q}}^2 \gg M_3^2) \end{cases}. \quad (8)$$

If  $M_{\text{unif}}$  is identified as the GUT scale, then  $c$  is typically in the range  $-1 \lesssim c \lesssim -\frac{1}{4}$ . Thus, this second contribution can either dramatically increase  $\epsilon_u$  or potentially cancel

much of it off, depending on their relative (model-dependent) signs. Perhaps more importantly, this contribution can still lead to large  $\epsilon_u$  even if the  $A$ -terms at the weak scale are small compared to the squark masses.

Now we return to Eq. (1). We can simplify it considerably by working in a basis in which  $\mathbf{Y}_\mathbf{U} = \mathbf{U}$  and  $\mathbf{Y}_\mathbf{D} = \mathbf{D}\mathbf{V}^{0\dagger}$  where  $\mathbf{V}^0$  is the CKM matrix at lowest-order (the meaning of this will be clear shortly) and  $\mathbf{U}$  and  $\mathbf{D}$  are both diagonal. Then

$$-\mathcal{L}_{eff} = \overline{D}_R \mathbf{D}\mathbf{V}^{0\dagger} Q_L H_d + \overline{D}_R \mathbf{D}\mathbf{V}^{0\dagger} [\epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{U}] Q_L H_u^* + h.c. \quad (9)$$

It is clear that in the absence of the  $\epsilon_u$  term, all pieces of the effective Lagrangian can be diagonalized in the same basis, preventing the appearance of flavor-changing neutral currents (FCNCs). It is the presence of the  $\epsilon_u \mathbf{U}^\dagger \mathbf{U}$  piece, however, that will prevent simultaneous diagonalization and generate some flavor-changing.

To see how this works, it is sufficient to keep only the Yukawa couplings of the third generation so that  $(\mathbf{U})_{ij} = y_t \delta_{i3} \delta_{j3}$  and  $(\mathbf{D})_{ij} = y_b \delta_{i3} \delta_{j3}$ . The flavor-conserving pieces of  $\mathcal{L}_{eff}$  then have the form

$$(1 + \epsilon_g) \mathbf{D}\mathbf{V}^{0\dagger} = (1 + \epsilon_g) y_b \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ V_{ub}^0 & V_{cb}^0 & V_{tb}^0 \end{pmatrix} \quad (10)$$

while the flavor-changing piece has the form

$$\epsilon_u \mathbf{D}\mathbf{V}^{0\dagger} \mathbf{U}^\dagger \mathbf{U} = \epsilon_u y_t^2 y_b \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & V_{tb}^0 \end{pmatrix}. \quad (11)$$

We can define a physical eigenbasis by rotating the  $d$ -component of  $Q_L$  by a new matrix  $\mathbf{V}$  defined by diagonalizing the mass matrix:

$$(\mathbf{V}^\dagger \mathcal{Y}^\dagger \mathcal{Y} \mathbf{V})_{ij} = \text{diag}(\overline{y}_d^2, \overline{y}_s^2, \overline{y}_b^2) \quad (12)$$

where the  $\overline{y}_i$  are defined to be the “physical” Yukawa couplings, *e.g.*,  $m_b = \overline{y}_b v_d$ ; and

$$\mathcal{Y} = \mathbf{D}\mathbf{V}^{0\dagger} [1 + \tan \beta (\epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{U})], \quad (13)$$

the  $\tan \beta$  coming from the vev of  $H_u$  which multiplies the loop-induced terms.  $\mathbf{V}$  can now be interpreted as the physical CKM matrix.

In the physical basis, the (3,3) element of the mass matrix gives us the corrected  $b$ -quark mass:

$$\overline{y}_b \simeq y_b [1 + (\epsilon_g + \epsilon_u y_t^2) \tan \beta]. \quad (14)$$

To get to this equation we used the fact that one finds no large (*i.e.*,  $\tan \beta$ -enhanced) corrections to  $V_{tb}$  [4], so that we can replace  $V_{tb}^0 \simeq V_{tb} \simeq 1$ .

The corrected CKM elements are the elements of  $\mathbf{V}$ . In particular,

$$V_{ub} \simeq V_{ub}^0 \left[ \frac{1 + \epsilon_g \tan \beta}{1 + (\epsilon_g + \epsilon_u y_t^2) \tan \beta} \right]. \quad (15)$$

The same form also holds for the corrected  $V_{cb}$ ,  $V_{td}$  and  $V_{ts}$ . Consistent with our earlier simplification that  $V_{tb}^0 \simeq 1$ , one finds that  $V_{tb}$  receives no correction. Note that Eqs. (14)–(15) present a coherent picture of the radiative corrections generated by SUSY-breaking: all of the diagrams represented in Fig. 1 contribute to a renormalization of the mass, but only the higgsino-mediated diagrams contribute a piece to the mass matrix which is not diagonal in the usual mass basis and which therefore generates FCNCs. Thus we see that  $V_{ub}$  reduces to  $V_{ub}^0$  in the limit that  $\epsilon_u = 0$ .

For  $\epsilon_u \neq 0$ , however, the rotation that diagonalized the mass matrix does not diagonalize the Yukawa couplings of the Higgs fields. Redefining  $D_L$  and  $D_R$  as the mass eigenstates, the effective Lagrangian for their couplings to the neutral Higgs fields is

$$-\mathcal{L}_{d,eff} = \overline{D}_R \mathbf{D} \mathbf{V}^{0\dagger} \mathbf{V} D_L H_d^0 + \overline{D}_R \mathbf{D} \mathbf{V}^{0\dagger} [\epsilon_g + \epsilon_u \mathbf{U}^\dagger \mathbf{U}] \mathbf{V} D_L H_u^{0*} + h.c. \quad (16)$$

Keeping only the flavor changing pieces, this simplifies after some algebra to

$$\mathcal{L}_{FCNC} = \frac{\overline{y}_b V_{tb}^*}{\sin \beta} \chi_{FC} [V_{td} \overline{b}_R d_L + V_{ts} \overline{b}_R s_L] (\cos \beta H_u^{0*} - \sin \beta H_d^0) + h.c. \quad (17)$$

with the quark fields in the physical/mass eigenbasis, and defining

$$\chi_{FC} = \frac{-\epsilon_u y_t^2 \tan \beta}{(1 + \epsilon_g \tan \beta)[1 + (\epsilon_g + \epsilon_u y_t^2) \tan \beta]} \quad (18)$$

to parameterize the amount of flavor-changing induced. Note also that we have expressed the coupling in terms of  $\overline{y}_b = m_b/v_d$  instead of the original, but unphysical,  $y_b$ .

The final step is to define the neutral Higgs mass eigenstates. These are defined as usual:

$$\begin{aligned} h^0 &= \sqrt{2}(\cos \alpha \operatorname{Re} H_u^0 - \sin \alpha \operatorname{Re} H_d^0), \\ A^0 &= \sqrt{2}(\cos \beta \operatorname{Im} H_u^0 + \sin \beta \operatorname{Im} H_d^0) \end{aligned} \quad (19)$$

and  $H^0$  orthogonal to  $h^0$ . Then the flavor-changing couplings between the Higgs mass states and the fermion mass states are:

$$\left. \begin{aligned} h^0 \overline{b}_R d_L : & \quad i \cos(\beta - \alpha) \\ H^0 \overline{b}_R d_L : & \quad i \sin(\beta - \alpha) \\ A^0 \overline{b}_R d_L : & \quad 1 \end{aligned} \right\} \times \frac{\overline{y}_b V_{td} V_{tb}^*}{\sqrt{2} \sin \beta} \chi_{FC} \quad (20)$$

A similar expression holds for the Higgs couplings to  $\overline{b}_R s_L$  with  $V_{td}$  replaced by  $V_{ts}$ . One non-trivial check of this result is to take the Higgs decoupling limit in which

$m_{A^0} \rightarrow \infty$ , driving  $\alpha \rightarrow \beta - \frac{\pi}{2}$ . There the  $h^0 \bar{b}_R d_L$  coupling goes to zero as it should in any single Higgs doublet model.

We will now consider two processes which constrain and/or provide a signal for the Higgs-mediated FCNCs:  $B^0 - \bar{B}^0$  mixing and the decay  $B^0 \rightarrow \mu^+ \mu^-$ . The case of  $B^0 - \bar{B}^0$  mixing is actually quite amusing.  $\Delta m_{B_d}$  is very well known and usually provides one of the tightest constraints on new sources of flavor-violation in the  $d$ -quark sector. And, in principle, mixing can be generated by single Higgs exchange. The leading order contribution of the 3 physical Higgs bosons to an effective operator  $\bar{b}_R^i d_L^i \bar{b}_R^j d_L^j$  ( $i, j$  are SU(3) indices) is proportional to the product of vertex factors and propagators given by:

$$\mathcal{F} \equiv \left[ \frac{\cos^2(\beta - \alpha)}{m_h^2} + \frac{\sin^2(\beta - \alpha)}{m_H^2} - \frac{1}{m_{A^0}^2} \right]. \quad (21)$$

However,  $\mathcal{F} = 0$  at lowest order. The existence of this zero is essentially an accidental cancellation coming from the special form of Eq. (17) and not an indication that the flavor-changing is illusory.

It is natural to ask whether this zero survives loop corrections, and one finds that it does not. However, the cost of adding another loop to the diagram is high and tends to suppress this new contribution too much to dominate the Standard Model contribution. We have considered in detail the largest non-zero contribution, which arises from top-stop induced vacuum polarization on the internal Higgs line. While these propagator corrections to the Higgs are known to be large [10], we find that the leading term (which is a correction to the  $H_u$  line) is suppressed by  $1/\tan^2 \beta$ . The next-leading term (a correction on the  $H_d$  line due to left-right stop mixing) is present but is not very large. All other radiative corrections we expect to be even smaller.

One can still derive a bound on  $m_A$  by demanding that the MSSM contribution to  $\Delta m_{B_d}$  is less than its observed value. Such a bound will depend sensitively on whether or not the two contributions to  $\epsilon_u$  from Eqs. (4) and (8) interfere constructively or destructively. Assuming all MSSM masses to be near 500 GeV and constructive interference, we find  $m_A \lesssim 100$  to 125 GeV for  $\tan \beta = 40$  to 60. Direct search constraints aside, it is known that models with such a light second Higgs doublet generally contribute far too much to  $b \rightarrow s \gamma$  and are therefore already ruled out [11]. Thus this new source of flavor-changing rules out a part of parameter space which is already known to be disfavored.

We now consider the rare decay  $B^0 \rightarrow \mu^+ \mu^-$ . This occurs via emission off the quark current of a single virtual Higgs boson which then decays leptonically. The largest leptonic flavor-changing branching fraction would clearly be to  $\tau^+ \tau^-$ . However, the branching fraction to  $\mu$ 's is only suppressed by  $(m_\mu/m_\tau)^2$  times a phase space factor, which is only about 1 part in 100. The current experimental limits on  $Br(B^0 \rightarrow \mu^+ \mu^-)$  are at the  $10^{-6}$  level, which means that the largest the branching ratio into  $\tau$ 's could be is about  $10^{-4}$ . Given the extreme difficulties encountered in trying to measure this decay experimentally, it is doubtful that the  $\tau$ -mode will ever

provide an interesting constraint or signal in and of itself. Thus we will concentrate on the  $\mu$ -channel.

The amplitude for the process  $B_{(d,s)}^0 \rightarrow \mu^+ \mu^-$  is given by:

$$\mathcal{A} = \eta_{QCD} \frac{\bar{y}_b y_\mu V_{t(d,s)} V_{tb}^*}{2 \sin \beta} \chi_{FC} \langle 0 | \bar{b}_R d_L | B_{(d,s)}^0 \rangle [\bar{\mu} (a_1 + a_2 \gamma^5) \mu] \quad (22)$$

where

$$\begin{aligned} a_1 &= \frac{\sin(\beta - \alpha) \cos \alpha}{m_H^2} - \frac{\cos(\beta - \alpha) \sin \alpha}{m_h^2}, \\ a_2 &= -\frac{\sin \beta}{m_A^2}. \end{aligned} \quad (23)$$

The partial width is then

$$\Gamma(B_{(d,s)}^0 \rightarrow \mu^+ \mu^-) = \frac{\eta_{QCD}^2}{128\pi} m_B^3 f_B^2 \bar{y}_b^2 y_\mu^2 |V_{t(d,s)}^* V_{tb}|^2 \chi_{FC}^2 (a_1^2 + a_2^2). \quad (24)$$

In the large  $m_A$ , large  $\tan \beta$  limit,  $a_1^2 + a_2^2 \simeq 2/m_A^4$ . The QCD correction is identical to the usual running of a quark mass operator, which in this case gives  $\eta_{QCD}$  between 1.4 and 1.6 for  $m_A$  between  $m_Z$  and 500 GeV. Experimentally,  $Br(B_{(d,s)}^0 \rightarrow \mu^+ \mu^-) < (6.8, 20) \times 10^{-7}$  at 90% confidence [12]. Thus  $\Gamma_{(d,s)} < (2.9, 8.7) \times 10^{-19}$  GeV. The factor of 3 in going from the  $B_d^0$  to the  $B_s^0$  limits is due to a factor of 3 suppression in the production cross-section of  $B_s^0$  compared to  $B_d^0$  at the Tevatron. However, theory predicts the partial width for  $B_s^0 \rightarrow \mu^+ \mu^-$  to be enhanced by  $(V_{ts}/V_{td})^2 \simeq 25$ . Thus one expects a signal in  $B_s^0$  decays before one is observed in  $B_d^0$ .

A few quick estimates can give us an impression of the importance of these new contributions. For nearly-degenerate MSSM particles at 500 GeV, one finds  $|\epsilon_g| \approx 1/80$  and  $|\epsilon_u| \approx (1/4)|\epsilon_g|$ , not including in  $\epsilon_u$  the contribution of Eq. (8). We derive a bound on  $m_A$  from the limit on  $B_s^0 \rightarrow \mu^+ \mu^-$  and using  $f_B = 180$  MeV and  $|V_{ts}| = 0.04$ . The bound depends sensitively on the signs of  $\epsilon_g$  and  $\epsilon_u$  as well as the size of the  $c$ -parameter of Eq. (8), which we take in the range  $-3/4 \leq c \leq 0$ . We also demand that  $y_b \leq y_t$  to avoid problems with perturbation theory and consistency with unification; this places an upper bound on  $\tan \beta$  as a function of  $\epsilon_g$ ,  $\epsilon_u$  and  $c$ . Varying over all of these, the strongest bounds are

$$m_A > (225, 175, 230, 215) \text{ GeV} \quad (25)$$

for  $\tan \beta = (29, 65, 38, 65)$ ,  $c = (-3/4, 0, 0, -3/4)$  and the signs of  $\{\epsilon_g, \epsilon_u\}$  being  $(--, ++, -+, ++)$  respectively.

Like the case of  $B^0 - \bar{B}^0$  mixing, we are finding ourselves in the range already constrained by  $b \rightarrow s\gamma$  and direct searches. However, unlike the mixing case where the MSSM contribution was typically smaller than the Standard Model prediction,

here we are still far above the Standard Model which predicts  $Br(B_{(d,s)}^0 \rightarrow \mu^+\mu^-) \simeq (1.5, 35) \times 10^{-10}$  [13]. Thus further experimental data can significantly improve the bounds on  $m_A$  or find a non-zero signal induced by supersymmetry.

So what is implied for Run II at the Tevatron? Assuming no change in their efficiencies and acceptances, CDF can in principle place a bound  $Br(B_s^0 \rightarrow \mu^+\mu^-) < 1 \times 10^{-7}$  given  $1 \text{ fb}^{-1}$  of data, a factor of 20 stronger than present. Thus the region probed in  $m_A$  will increase by  $20^{1/4} \simeq 2$ :

$$m_A > (475, 365, 490, 450) \text{ GeV} \quad (26)$$

for the same sets of inputs as previously. After collecting  $5 \text{ fb}^{-1}$  these masses increase by another 50%, up to 725 GeV. Finally, if the proposed “Run III” of the Tevatron with  $30 \text{ fb}^{-1}$  were to occur, masses of  $A^0$  all the way to 1 TeV could be studied. This could be a very important signal for supersymmetry since *this source of flavor-changing does not decouple as  $M_{\text{SUSY}} \rightarrow \infty$*  so long as  $m_A$  does not also get very heavy. That is to say, the bound on  $m_A$  is roughly independent of  $M_{\text{SUSY}}$ . Therefore supersymmetric spectra in the multi-hundred GeV to TeV range may be probed at the Tevatron through rare  $B$ -decays even when direct production of supersymmetry (including the second Higgs doublet) cannot be observed. Since the precise predictions for  $Br(B_s^0 \rightarrow \mu^+\mu^-)$  are highly dependent on the individual model, these estimates should only be taken as indicative. Further work will be forthcoming [8].

It is also possible to look for new sources of flavor-changing in inclusive semileptonic decays  $B \rightarrow X_s \mu^+ \mu^-$ . The width for this process can be extracted from Eq. (24) by replacement of  $f_B$  with  $m_B$  and dividing by  $192\pi^2$  for the 3-body phase space. The rate is thus a factor of 10 smaller than for  $B_s \rightarrow \mu^+ \mu^-$ . Comparing to current bounds [14] yields constraints on  $m_A$  that are weaker by a factor of 1.8 than the bounds from the purely leptonic mode. The ability of future experiments to extract information from this mode will be discussed in [8].

Finally, we find it noteworthy that the largest signals tend to occur for  $\epsilon_g < 0$  and intermediate values of  $\tan \beta$ . In minimal GUT models, one expects unification of the  $b$ - and  $\tau$ -Yukawa couplings. But it is well-known that this unification fails over most of the parameter space of the MSSM and generally necessitates the use of the HRS corrections to bring the Yukawas back into agreement. Typically one requires  $(\epsilon_g + y_t^2 \epsilon_u) \tan \beta \approx -0.2$  [3, 15] which in turn means that  $\epsilon_g < 0$ . This provides an argument for believing that the signal might lie in the observable range, as well as providing another test of Yukawa unification (beyond those discussed in [6] for flavor-conserving processes).

In summary, we have found that neutral Higgs bosons are capable of mediating flavor-changing interactions within the MSSM. This result is generic and does not rely on assumptions about sparticle mass non-universality which are usually required in order to get FCNCs. These interactions are enhanced at large  $\tan \beta$  and are in the range that will be experimentally probed in the near future.



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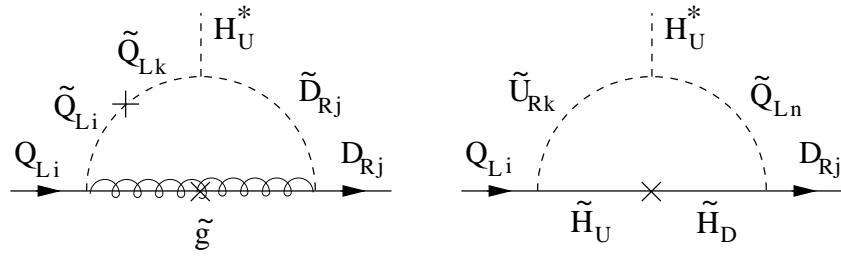


Figure 1: Leading contributions to  $\epsilon_g$  and  $\epsilon_u$ . Indices  $i, j, k, n$  label flavors.